

The frozen (inactive) disk in Sgr A*: freezing the accretion of the hot gas too?

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The black hole (BH) in our Galactic Center (GC) is extremely underluminous for the amount of hot gas available for the BH consumption. Theoretical understanding of this fact rests on a likely but not entirely certain assumption that the electrons in the accreting gas are much cooler than the protons. In this case the hot gas as a whole is too hot to accrete, and is too tenuous to radiate away its gravitational energy. Here we propose a drastically different picture of the accretion process in Sgr A* not based on the unchecked two-temperature assumption. Namely, we argue that there should exist a very cold *inactive* disk – a remnant of a past stronger accretion activity in Sgr A*. Such a disk would be a very efficient *cooling surface* for the hot flow. We show that under certain conditions the cooling due to thermal conduction cannot be balanced by the viscous heating in the hot flow. Along with the heat, the hot flow loses its viscosity and thus ability to accrete. It settles (condenses) onto the cold disk slightly inside of the circularization radius. If the latter is very large, then the liberated energy, and the luminosity emitted, is orders of magnitude less than naively expected. We build a simple analytical model for this flow and calculate the expected spectra that appear to be in a very reasonable agreement with observations. Strong additional support for the presence of the inactive disk comes from the recent observations of X-ray flares in Sgr A*. The properties of these flares are very similar to those produced by stars passing through a cold disk.

1 Introduction

Currently, the pitiful luminosity of Sgr A* is most commonly explained in the framework of Non-Radiative Accretion Flows (NRAF), a generalization of solutions discussed in greatest detail by Narayan & Yi (1994; NY94 hereafter). These solutions are valid when the rate at which the protons pass their gravitational energy to the electrons is not significantly higher than that due to Coulomb interactions only. This assumption is unfortunately prohibitively difficult to test (see references in Narayan 2002). However, provided it is valid, the electrons radiate only a tiny fraction of the total energy (NY94, Narayan et al. 1995) in stark contrast to the standard disks (Shakura & Sunyaev 1973). The second important feature of NRAFs was pointed out by Blandford & Begelman (1999; BB99 hereafter) who showed that these flows should produce powerful thermally driven winds (see also Quataert, these proceedings).

In this paper we would like to point out a likely and essential element of the accretion picture in Sgr A* that has so far escaped (except for Falcke & Melia 1997; FM97) the attention it deserves. Sgr A* is believed to be closely related to the Low Luminosity AGN (LLAGN; e.g. Ho 1999). Most if not all of these sources seem to have cold *neutral and inactive* disks that often can be seen only through water maser emission (e.g. Miyoshi et al. 1995) arising in a range of radii where gas temperature is 200 – 1000 K. Continuous SED spectra of these objects also support existence of cold disks (Quataert et al. 1999). Because of the extremely low ionization level of these cold disks, the disk viscosity may be nearly zero (e.g. see

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Menou & Quataert 2001). As such these “frozen” disks are not *accretion* disks (e.g. Siemiginowska et al. 1996). Falcke & Melia (1997) studied the evolution of such a disk on very long time scales in Sgr A* and concluded that unless the stellar wind (which is the current source of gas for the hot flow) has a very large angular momentum, it would have violated the near infra-red limits. In §4 we argue that due to some more recent data (e.g. Genzel 2000), the wind *does have* a substantial angular momentum and hence the problem needs to be re-considered.

Consider now the implications of an inactive disk presence for the flow of the hot gas in Sgr A*. The cool disk is razor-thin but is much more massive than the hot flow (Nayakshin et al. 2003). For temperatures as high as 10^7 K, thermal conduction is usually very important. For the conditions at hand, the time to cool off via thermal conduction, t_{cond} , is much shorter than that by radiation. Moreover, this time can even be shorter than the flow viscous time. In this case the hot flow gets “frozen” by thermal conduction before it can accrete onto the BH. As the flow looses its thermal energy, it also looses its vertical pressure support, and therefore it has to settle down (condensate) onto the inactive disk. Becoming a part of the inactive disk, the gas looses its viscosity and its ability to accrete, and can stay in the same location essentially indefinitely (the cold disk viscous time is as long as $10^6 - 10^8$ years at large radii). The accretion process in this picture would be “delayed” because the gas is currently piled up in the cold disk. The expected bolometric luminosity of this “frozen flow” is roughly $L_{\text{bol}} \sim 0.1 \dot{M} c^2 (\text{few } R_g/R_c)$ where R_c is the circularization radius of the hot flow and $R_g = 2GM_{\text{BH}}/c^2$ is gravitational radius. If $R_c/R_g \gtrsim 10^4$, then the low luminosity of Sgr A* could be understood naturally: the hot gas does not penetrate very deep into the BH potential well, thus gaining little energy as it settles onto the inactive disk.

During the workshop we presented 2D hydrodynamical simulations of the hot flow above a cold inactive disk. Using an insight from our simulations, we have recently discovered a simple idealized analytical solution that provides a much easier understanding of the numerical results. We present this solution below. The spectra resulting from this condensing flow appear to agree quite well with the observational constraints if condensation radius is $R_c \gtrsim 3 \times 10^4 R_g$, the disk is highly inclined, i.e. $i \gtrsim 75^\circ$, and the accretion rate in the hot flow is at its “nominal” value, $\dot{M}_0 \sim 3 \times 10^{-6} M_\odot/\text{year}$.

In addition, during the meeting we realized that the observed X-ray flares may well be due to stars passing through the inactive disk. In Nayakshin & Sunyaev (2003; NS03 hereafter) and in Nayakshin et al. (2003) we calculated the expected rate of flares, duration of a typical flare, X-ray spectra, luminosities, plus flare radio and NIR luminosities. All of these quantities closely resemble the observational picture reported by Baganoff et al. (2001) & Baganoff et al. (2003). It appears to us that both quiescent and flare spectra of Sgr A* can be explained if we accept existence of an inactive frozen disk.

2 A simple analytical model

A disk with temperature $T_d \lesssim 100$ K and with outer radius $R_d \lesssim \text{few} \times 10^4 R_g$ could be very hard to detect in Sgr A* with any current telescopes (NS03) if the disk is also highly inclined. Since the disk is razor-thin, its viscous time scale is very large ($\gtrsim 10^5$ years), and we can consider its structure as given on shorter time scales. We thus introduce the disk only through the boundary conditions for the hot flow.

Two-component accretion flows are too complex to be studied analytically except for very simplified special cases (which can however be most insightful – see BB99 for an example). Therefore, we will restrict ourselves to vertically averaged equations for a Keplerian hot flow. We consider radii $R < R_c$ where R_c is the circularization radius of the hot gas. For the low densities concerned, the radiative cooling of the hot gas is negligible, while the thermal conduction is at its “best” – at the maximum or the “saturated” value. Namely, the heat flux is given by $F_{\text{sat}} = 5\phi P c_s$, where $\phi \leq 1$ is the saturation parameter (Cowie & McKee 1977), and P is the vertically averaged pressure, $P = \rho c_s^2$ (ρ and c_s are the gas density and the isothermal sound speed, respectively). Note that if $\phi \sim 1$ then the energy flow due to thermal conduction is effectively supersonic. This is because the flux is carried by the electrons whose thermal speed is much higher than that of the protons in a one-temperature plasma (for more on this see Cowie & McKee 1977).

On the other hand, the flow of energy in the radial direction will be proportional to the radial velocity that is normally much smaller than the sound speed (see below). Thus the energy losses due to thermal conduction in the limit of large ϕ are much larger than the rate at which the energy can be gained by sinking in the BH potential well. That is why condensation should ensue in this situation.

Our equations are best understood through comparison with those of the well known ADAF solution (NY94; see their eqs. 1-4). The stationary mass conservation equation takes into account exchange of mass in the vertical direction:

$$\frac{\partial}{\partial R} [RH\rho v_R] = -\rho v_z R, \quad (1)$$

where v_R is the radial flow velocity ($v_R > 0$ for accretion), ρv_z is the mass flow density for condensation ($v_z < 0$) or evaporation ($v_z > 0$). The vertical scale height H is introduced through the hydrostatic balance as $H = c_s/\Omega$ where Ω is the angular velocity. Since the latter is assumed to be Keplerian, the radial momentum equation (eq. 2 in NY94) is trivially satisfied. Since the cold disk is also Keplerian, there is no exchange of specific angular momentum between the two flows and the angular momentum conservation equation (3 in NY94) is unaltered. With $\Omega = \Omega_K$ we get

$$\dot{M} = 4\pi RH\rho v_R = \frac{12\pi\alpha}{R\Omega_K} \frac{\partial}{\partial R} [\rho c_s^2 R^2 H], \quad (2)$$

where α is viscosity parameter (Shakura & Sunyaev 1973). The entropy equation should include (in addition to the usual terms) the thermal conduction flux, F_{sat} , and the hydrodynamical flux of energy in the vertical direction. To derive this equation, we follow the formalism of Meyer & Meyer-Hofmeister (1994; see their equation 8), with the following exceptions. We neglect winds here because we are interested in cooler, condensing solutions. Thus their side-way term (their eq. 5) is not included. In addition, the radial entropy flow term (the “advective cooling”; NY94) is designated Q_{adv} . Following NY94 we set $Q_{\text{adv}} = f_{\text{adv}}Q_+$, where $f_{\text{adv}} \leq 1$ is a parameter. Further, for subsonic flows $v_z^2 \ll c_s^2$, and we obtain

$$Q_+ - F_{\text{sat}} - \rho v_z \left[\frac{5}{2}c_s^2 + \frac{GM_{\text{BH}}}{2R} \left(\frac{H}{R} \right)^2 \right] - Q_{\text{adv}} = 0, \quad (3)$$

where $Q_+ = (9/2)\alpha\rho c_s^3$ is the viscous heating rate. The factor of $5/2$ in eq. (3) is $\gamma/(\gamma - 1)$ for the $\gamma = 5/3$ gas that we consider here.

After some simple algebra and using $GM_{\text{BH}}H^2/2R^3 = c_s^2/2$, we find that the condensation velocity is

$$v_z = - \left(\frac{5}{3}\phi - \frac{3}{2}\alpha(1 - f_{\text{adv}}) \right) c_s = -bc_s, \quad (4)$$

where b is introduced for convenience. The equation (4) is crucial for the rest of the paper. If thermal conduction is vigorous, i.e. $b > 0$, then the hot flow is condensing onto the cold disk. In the opposite case of a small ϕ and a “large” α , $b < 0$, viscous heating prevails. Thermal conduction then serves to evaporate the inactive disk. We should also note that the regime of large α (~ 0.3) was already studied by F. Meyer and collaborators in many papers. Their solutions are for higher accretion rates and therefore they are in the non-saturated regime, which is roughly speaking equivalent to the $\phi \ll 1$ case. As they found, the evaporation is a very strong function of α ($\propto \alpha^3$). Our results (eq. 4 in particular) are thus in a complete agreement with that of Meyer & colleagues, and we essentially extend their work on the case $\phi \gg \alpha$.

Inserting now $v_z = -bc_s$ into equation (1) and also substituting v_R on its value found from equation (2), we arrive at a second order differential equation that contains two variables, ρ and c_s :

$$b\rho Rc_s = 3\alpha \frac{\partial}{\partial R} \left\{ \frac{1}{R\Omega_K} \frac{\partial}{\partial R} \left[\rho c_s^3 \frac{R^2}{\Omega_K} \right] \right\}. \quad (5)$$

This equation cannot be solved in a general case. By introducing $v_z \neq 0$ we added an extra variable to the accretion flow equations, and the number of independent equations is now smaller than the number of unknowns. (In particular, both v_z and c_s are to be found from the single energy equation 3). This situation is well known in analytical ADIOS wind solutions (BB99). In the latter case one has to introduce three free parameters that describe the mass, energy and angular momentum carried away by the wind.

The most natural way to proceed here is to suggest that the temperature is a power-law function of radius. For example, for ADAF, $T(R) \propto R^{-1}$ in a broad range of radii. On the other hand, thermal conduction tends to smooth out temperature gradients (for example within supernova remnants), and hence in the other extreme $T(R) \simeq \text{const}$ (we in fact observed this nearly constant $T(R)$ in our numerical simulations). Thus, $c_s = c_0(R_c/R)^\lambda$, where $0 \leq \lambda \leq 1/2$ and c_0 is the sound speed at R_c . If we define $u \equiv \sqrt{R/R_c}$, then equation (5) can be re-written as

$$\rho u^{3-2\lambda} = \frac{3\alpha}{4b} \frac{R_c c_0^2}{GM_{\text{BH}}} \frac{\partial^2}{\partial u^2} [\rho u^{7-6\lambda}] . \quad (6)$$

Finally, defining $\tilde{\rho} \equiv \rho u^{7-6\lambda}$, we obtain

$$\frac{\partial^2 \tilde{\rho}}{\partial u^2} = \frac{1}{l_c^2} \frac{\tilde{\rho}}{u^{4(1-\lambda)}} . \quad (7)$$

Here l_c is the dimensionless “condensation length”:

$$l_c^2 \equiv \frac{3\alpha}{4b} \frac{c_0^2 R_c}{GM_{\text{BH}}} . \quad (8)$$

If there were no thermal conduction losses, the internal energy of the hot gas would be of order its gravitational energy. If the inactive disk extends to $R > R_c$, the thermal conduction will reduce the gas thermal energy. Hence we expect that $c_0^2 R_c / GM_{\text{BH}} \leq 1$.

We are most interested in the case $\phi \gg \alpha$ and therefore we shall only explore the $\lambda = 0$ case below.¹ In addition, with $\alpha \ll \phi$, condensation length is small, i.e. $l_c^2 < \alpha/b \sim \alpha/\phi \ll 1$. This circumstance facilitates finding an approximate analytical solution for equation (6):

$$\tilde{\rho} = \text{const} \exp \left[-\frac{1}{l_c u} \right] . \quad (9)$$

Indeed, $d^2 \tilde{\rho} / du^2 = \tilde{\rho} (-2/l_c u^3 + 1/l_c^2 u^4) \simeq \tilde{\rho} / l_c^2 u^4$ due to the fact that $1/l_c u \gg 1$. Let us now explicitly write down the results from this approximate solution:

$$\rho(R) = \rho_c \left(\frac{R_c}{R} \right)^{7/2} \exp \left[-\frac{1}{l_c} \left(\sqrt{\frac{R_c}{R}} - 1 \right) \right] , \quad (10)$$

$$v_R = \frac{3\alpha}{2l_c} \frac{c_0^2}{R_c \Omega_K(R_c)} = \sqrt{\frac{\alpha b}{3}} c_0 = \text{const} \ll c_0 , \quad (11)$$

$$\dot{M}(R) = \dot{M}_0 \frac{R_c}{R} \exp \left[-\frac{1}{l_c} \left(\sqrt{\frac{R_c}{R}} - 1 \right) \right] . \quad (12)$$

Here ρ_c is gas density at R_c , and \dot{M}_0 is the accretion rate at that point. Note that equation (11) shows that the radial velocity is constant and is substantially smaller than the sound speed. Further, for $l_c < 1/2$ (recall that we assumed $l_c \ll 1$), the accretion rate increases with R , as it should for a condensing flow.

¹ Nevertheless we note that (i) approximate solutions may be obtained for a general value of λ , and (ii) for $\lambda = 1/2$ there is a scale-free solution.

3 Sample results

To illustrate the analytical solution presented above, we plot the gas density and the mass accretion rate profiles (equations 10 & 12) in Figure 1 for $\phi = 0.2$ and two different values of α . These are chosen to represent the two opposite extremes. In the case of larger α , condensation length is $l_c = 0.38$ and this is barely satisfies the condition $l_c \ll 1$ under which our approximate solution is valid. The accretion rate changes slowly with radius in this case, meaning that condensation is relatively slow. In the case of very small α the hot flow collapses onto the disk far quicker: the mass flow is reduced by ~ 100 times already at $R \simeq R_c/2$. This is easily understood by noting that $v_z/v_R = \sqrt{3b/\alpha} \simeq 4$ for the larger α value, whereas for $\alpha = 0.004$ this ratio is $v_z/v_R \simeq 16$. The large value of the ratio v_z/v_R in our solutions justifies the statement that we made in the introduction: “condensation time”, $t_{\text{cond}} = H/v_z$, is much shorter than the viscous time, $t_{\text{visc}} = R/v_R$ for our solutions.

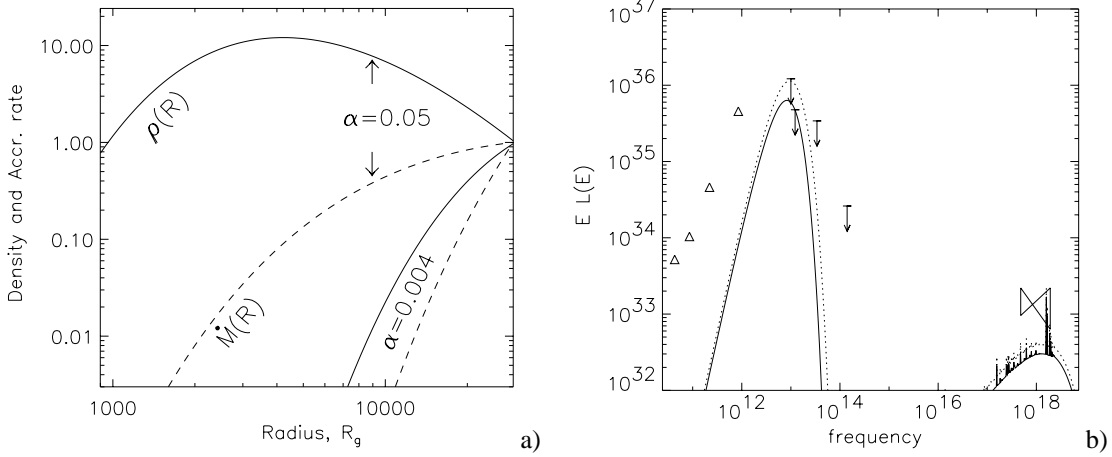


Fig. 1 (a) – Radial density (solid) and accretion rate (dashed) profiles for $\phi = 0.2$, $R_c = 3 \times 10^4 R_g$, and two values of viscosity parameter $\alpha = 0.05$ (upper curves) and 0.004 (lower curves). (b) – Spectra corresponding to the radial profiles shown in (a). The dotted and solid curves are for $\alpha = 0.05$ and $\alpha = 0.004$, respectively. For both cases inclination angle $i = 75^\circ$ and accretion rate $\dot{M}_0 = 3 \times 10^{-6} M_\odot \text{ year}^{-1}$. The time-averaged spectrum of Sgr A* is shown with triangles (radio data; detections), upper limits (infra-red) and the bow-tie (X-rays; Chandra data). The lower frequency data are from references given in Melia & Falcke (2001), except for the $2.2 \mu\text{m}$ point that is from Hornstein et al. (2003), and for Chandra data that are from Baganoff et al. (2003). For the X-ray part of the data the spectrum should only be lower than the Chandra observations because the latter are probably dominated by the emission of hot gas at $R \sim R_B > R_c$, which is not included in the model.

It is useful to define the “apparent” bolometric efficiency coefficient of the frozen accretion flows, ε_{bol} in the standard way:

$$\varepsilon_{\text{bol}} \equiv \frac{L_{\text{bol}}}{\dot{M}_0 c^2}, \quad (13)$$

where L_{bol} is the bolometric luminosity of the source. One should recall that for standard accretion flows, $\varepsilon_{\text{bol}} \sim 0.1$ for a non-rotating BH. The energy deposition into the cold disk is given by the second and the third terms in equation (3). Integrating this expression over the disk surface area from $3R_g$ to R_c , one can arrive at $\varepsilon_{\text{bol}} = 1.5(1 + \zeta)[1 + 2l_c + 2l_c^2](R_g/R_c)$, where $\zeta \equiv (1 - 0.9\alpha/\phi)^{-1}$. For $\phi \gg \alpha$ we have $\zeta \simeq 1$ and $l_c \ll 1$, so $\varepsilon_{\text{bol}} \simeq 3R_g/R_c$. This expression is of course only applies at $R_c \gg 3R_g$ because we neglected any relativistic corrections to the gravitational potential. In addition, the derived expression is only a rough estimate since we assumed vertically averaged equations which are clearly inaccurate for $H \sim R$ (which in our simple model occurs at $R = R_c$). One can in fact show that realistically ε_{bol}

should be smaller by a factor of at least few. The point here is that our simplistic solution is not bound at $R = R_c$ because its thermal energy exceeds gravitational energy at that point. So either a wind takes away the excess energy (as BB99 argued for ADAF solutions) or more likely this simply means that we over-estimated the gas temperature at R_c . We should also explicitly insert the disk inclination angle into the definition of the apparent efficiency since the observed luminosity of the thin disk is approximately proportional to $\cos i$. Variation of the predicted spectrum with R_c is shown in Figure (3). Summarizing, the apparent bolometric efficiency of the freezing flow is

$$\varepsilon_{\text{bol}} \lesssim \cos i \frac{R_g}{R_c} = 10^{-4} \cos i \frac{10^4 R_g}{R_c} \quad (14)$$

In a similar spirit, the X-ray radiative efficiency can be defined. Approximately, $\varepsilon_x \sim (R_g/R_c) (t_{\text{cond}}/t_{\text{cool}})$, where t_{cool} is the radiative cooling time for the hot flow. For Sgr A*, ratio $t_{\text{cond}}/t_{\text{cool}}$ is extremely small and this is why X-ray contribution to the radiative output of Sgr A* is so small.

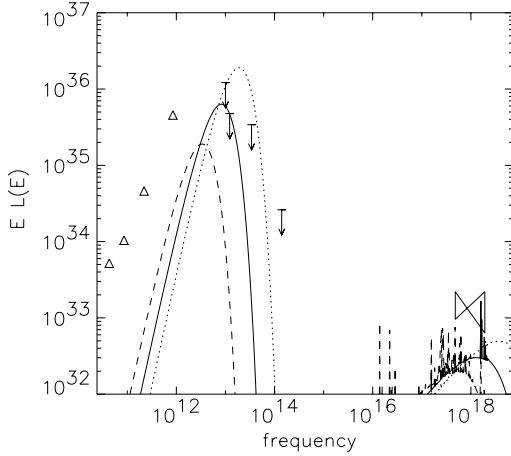


Fig. 2 Spectrum emitted by the condensing flow for inclination angle $i = 75^\circ$, accretion rate $\dot{M}_0 = 3 \times 10^{-6} M_\odot$ and three values of circularization radius R_c : $10^4 R_g$ (dotted), $3 \times 10^4 R_g$ (solid), and $10^5 R_g$ (dashed), respectively. As expected, the bolometric luminosity increases as R_c decreases because the hot gas gains more gravitational energy for larger R_c .

4 Discussion

We have suggested here that there exist an inactive disk around Sgr A*, a remnant of a past powerful accretion (and probably star forming) activity. The disk may be quite light compared with both the BH and the star cluster (Nayakshin et al. 2003), yet it easily out-weights the $\sim 10^{-3} M_\odot$ of the hot gas present in the region interior to the Bondi radius. The disk then serves as a very efficient cooling surface for the hot flow. The flow essentially gets *frozen* (stopped), and its energy is radiated as thermal emission at frequencies much below the X-ray band. The X-ray emitting flow thus simply disappears from “the radar screen”. This in our opinion may be the explanation of the exceptional *apparent* X-ray radiative inefficiency of Sgr A*. Further, since the hot flow does not penetrate very deep into the BH potential well, its total bolometric luminosity is \sim few R_g/R_c times smaller than that expected if the gas made it all the way into the BH and was radiatively efficient. Thus the flow also *appears* radiatively inefficient in the bolometric sense. As we explained in the paper, we believe that the accretion of the winds from the hot stars is simply delayed in time and it is by no means radiatively inefficient in the long run.

Falcke & Melia (1997; FM97) have assumed the hot wind infall as given and studied viscous evolution of the “fossil” disk on long time scales, whereas we concentrate on much shorter time scales on which the structure of the disk does not change (i.e. $t < t_{\text{visc}} \sim 10^6 \alpha^{-1}$ years for $T_{\text{disk}} = 100$ K and $R = 10^4 R_g$). Our study is therefore complimentary to that of FM97.

Our results concerning the conditions under which the wind–disk (or the hot flow–disk) interactions will not violate the tight NIR limits are quite similar to that of FM97. In particular, FM97 note that “the Bondi–Hoyle wind must be accreting with a very high specific angular momentum to prevent it from circularizing in the inner disk region where its impact would be most noticeable”. We find that the circularization radius should be $\gtrsim 3 \times 10^4 R_g$, implying a very large angular momentum indeed. Further, we suggested that the disk and the hot flow angular momenta are at least approximately aligned or else there would be a substantial heating due to friction between the two, a heating not included in our analysis. It remains to be seen whether results will be qualitatively similar if the disk and the hot flow rotation axes are misaligned.

FM97 considered a “large” value of R_c being rather unlikely. We however note that according to recent data, the stars from which the hot wind originates appear to be on tangential orbits counter-rotating the Galactic rotation and are \sim few arcsecond \sim few $\times 10^5 R_g$ off Sgr A* (e.g. Genzel 2000). There is thus no deficit of angular momentum at these distances. Finally, we only studied here the region of the flow interior to R_c . However the exchange of the angular momentum between the hot flow and the disk should take place at $R > R_c$ where the hot gas is sub-Keplerian. This should enrich the hot flow with the angular momentum directed as that of the the disk and hence the circularization of the hot flow should occur even if it had zero angular momentum at infinity. Thus the requirement of a large angular momentum in the wind may be relaxed, although this effect remains to be quantified with future calculations.

5 Conclusion

In this paper we suggested that there exists a very cold inactive disk in Sgr A*, and that its role in the accretion picture is significant. While the hot gas is very tenuous and cannot radiate its energy away, it can easily transfer its energy into the cold disk via thermal conduction. The cold disk is much denser and much more massive than the hot flow and can serve as a very powerful freezer (or radiator) for the hot flow. As the hot flow loses its energy, it also loses its viscosity and “sticks” to the cold disk. The accretion flow is thus quenched by this seemingly “non-radiative” cooling. One can easily check that neither internal viscous dissipation nor the heat input from the hot flow in Sgr A* are sufficient to overcome the radiative cooling and restart the accretion in the inactive disk. It appears that only arrival of a new large supply of low angular momentum material could revive the inactive disk in Sgr A* now.

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